Advances in Old–Fashioned Heterotic String Model Building

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I review findings of various research groups regarding perturbative heterotic string model building in the last 12 months. Attention is given to recent studies of extra $U(1)$'s and local discrete symmetries (LDS's) in generic string models. Issues covered include the role of $U(1)$'s and LDS's in limiting proton decay, developments in classification of models containing anomalous $U(1)$, and possible complications resulting from kinetic mixing between observable and hidden sector $U(1)$'s. Additionally, recent string-derived and string-inspired models are briefly reviewed.

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1. Perturbative String Model Building Circa 1996-1997

From one point of view string theory has undergone a profound paradigm shift in the last few years, a shift arguably so revolutionary that the term "string theory" will most likely be replaced by some appellation (e.g. "M"-theory or "F"-theory) for the current vaguely understood underlying theory. However, from a differing viewpoint, things have not significantly changed in some aspects. The new understanding of string dualities, connecting the various string "theories" formerly thought to be independent, does not appear to significantly effect physics in the weak coupling limit of heterotic string model building (HSMB). Thus, "old fashioned" HSMB is alive and well. Weakly-coupled heterotic strings can still play an important role in model building. This talk is a review of advancements made in HSMB over the course of the last year.

Section two of the talk discusses some generic features of perturbative heterotic string models (HSMs) whose presence has no field-theoretic explanation. In particular, the role of extra local abelian and local discrete symmetries (LDS's) in HSMs are reviewed. In the (quasi-)realistic HSMs to date, such symmetries play an important role in constraining proton decay (and flavor changing neutral currents), while simultaniously allow-ing light neutrino masses [1,7]. Next, progress in classification of models containing anomalous $U(1)$ are reviewed [8], including determination of anomalous $U(1)$ constraints in orbifolds. Section two concludes with discussion of HSMB complications resulting from possible kinetic mixing between observable and hidden sector $U(1)$'s [9]. Section three briefly examines recent progress in actual HSMB.

2. General Aspects of Model Building

The gauge structure of the more phenomenologically realistic four-dimensional HSMs follows the general form,

$$\{SU(3)_c \times SU(2)_l \times U(1)_Y \in G$$
$$\times \sum_i U(1)_i \times \sum_j \Delta_j\}_{\mathrm{obs}} \times \qquad (1)$$
$$\{\tilde{G}_{\mathrm{NA}} \times \sum_k \tilde{U}(1)_k \times \sum_l \tilde{\Delta}_l\}_{\mathrm{hid}},$$

where Δ_j denote LDS's. Minimally embedding the SM $SU(3)_c \times SU(2)_l$ symmetries through free fermionic realization [12] corresponds to realizing $SU(3)_c \times U(1)_c \in SO(6)$ from charges of three complex fermions and $SU(2)_l \times U(1)_l \in SO(4)$ from those of two complex fermions. This embedding fixes the SM hypercharge to be

$$U(1)_Y = \frac{1}{2}U(1)_l + \frac{1}{3}U(1)_c$$

$$= U(1)_{I_{3R}} + \frac{1}{2}U(1)_{B-L}, \qquad (2)$$

Non-minimal embeddings involve larger numbers of fermions and can offer degrees of freedom in the hypercharge realization. Differing hypercharge definitions can significantly alter the phenomenology of the effective field theory.

The separation between the observable and hidden sectors is generally not absolute. Indeed, most often some of the hidden sector states also carry observable $U(1)_i$ charges.[1] Charges carried by both hidden and observable sector states are often said to belong to the "shadow sector," Although states with both shadow and hidden sector charge are generally present, quasi-realistic phenomenology requires that they have masses $m_{\text{shadow}} \geq m_{\text{susy}}$.

In the conventional supergravity (SUGRA) senario, supersymmetry (SUSY) is broken in the hidden sector at a high scale

$$M_{\text{susy}}^{\text{hid}} \sim \sqrt{M_Z M_{\text{Pl}}} \sim 10^{10 \text{ to } 12} \text{ GeV}. \qquad (3)$$

Gravitational interactions (and possibly Planck-scale-suppressed non-renormalizable superpotential terms) then set the observable sector SUSY-breaking mass scale,

$$M_{\text{susy}}^{\text{obs}} \sim (M_{\text{susy}}^{\text{hid}})^2/M_{\text{Pl}} \sim M_Z. \qquad (4)$$

In the less conventional gauge-mediated process the hidden sector SUSY-breaking scale $M_{\text{susy}}^{\text{hid}}$ is typically around 100 TeV. SUSY-breaking is carried to the observable sector via loop-factor suppression factors.

Some interesting papers focusing on various roles of the ubiquitous stringy $U(1)$'s have appeared in the last year. In the following subsections I will discuss four issues in these papers.

2.1. $U(1)$ Beyond Those Within E$_6$

Recently Pati has presented evidence supporting the stringy origin of MSSM matter. Ref. [1] demonstrates that MSSM states in any phenomenologically viable model must carry at least one local $U(1)$ or LDS charge not originating from within the same E$_6$ (or subgroup thereof)

[2] that $SU(3)_c \times SU(2)_l \times U(1)_Y$ (henceforth denoted by (321)) may be embedded into. Extra $U(1)$'s are required for suppression of unsafe color-triplet mediated and/or gravity-linked proton decay operators generically arising in SUSY (grand)-unification. The necessity of an additional $U(1)$ becomes even more apparent when seesaw-type neutrino masses are desired. While extra $U(1)$ may be invoked ad hoc from the field-theoretic approach, the appearance of (several) additional $U(1)$'s is a natural outcome in most string models.

In MSSMs, the dangerous proton decay operators arise from baryon and lepton number violating superpotential terms of the form

$$\begin{aligned} W = & [\eta_1 U^c D^c D^c + \eta_2 QLD^c + \eta_3 LLE^c] \\ + & [\lambda_1 QQQL + \lambda_2 U^c U^c D^c E^c \qquad (5) \\ + & \lambda_3 LLH_2 H_2]/M, \end{aligned}$$

where SU(2), SU(3), and generational indices are suppressed. M is a characteristic mass scale, assumed herein to be $M_{\text{string}} \approx g_s \times 5.5 \times 10^{17}$ GeV, where the string coupling $g_s \approx 0.7$.

η_i and λ_j represent terms of generic order and can contain built-in suppression factors $(\langle\phi\rangle/M)^n$, where $\langle\phi\rangle$ is a non-Abelian singlet state VEV.[2] Proton decay limits imply $\eta_1\eta_2 \lesssim 10^{-24}$ and $\lambda_j/M(\text{GeV}) \lesssim 10^{-25}$ for $\Delta(B-L) = 0$ decays.

While dangerous dimension four ($d = 4$) operators from renormalizable η_i terms in (5) can be avoided in generic SUSY (GUT) models either by an ad hoc discrete symmetry like $R = (-1)^{3(B-L)}$ or by a local $U(1)_{B-L} \in SO(10)$ symmetry, neither symmetry can prevent similar effective $d = 4$ operators from non-renormalizable η_i terms containing factors of $(B-L)$ violating VEVs. Neither are the $d = 5$ operators associated with the $(B-L)$-conserving, non-renormalizable λ_1 and λ_2 terms forbidden. These latter terms are expected in any theory associated with gravity unless forbidden by a new symmetry. Additionally, in SUSY GUT models, even when sufficient sup-

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[1]Furthermore, three-generation free fermionic models exist that contain exotic states carrying both observable and hidden non-Abelian charges [13].

[2]$\langle\phi\rangle$ could, in general, be replaced by a singlet product of non-Abelian fields, such as a condensate, $\langle T\overline{T}\rangle$, of two hidden sector vector-like fields. In this case, additional factors of $1/M$ are then implied.

pression of λ_j is possible, induced $d = 5$ operators from Higgsino triplet exchange must be prevented by doublet-triplet splitting. These operators can be avoided by fine-tuning in SU(5) models or by suitable choices of large numbers of and/or high-dimensionality for multiplets, as is often the method in SO(10) models. However, such solutions do not seem compelling, especially when stringy constraints on allowed representations (reps) are considered [3]. (See footnote 4.)

In addition, the standard seesaw mechanism assigning light masses to the (primarily) ν and heavy masses to the (primarily) ν^c requires breaking $(B-L)$ at a high intermediate scale. However, violating $(B-L)$ by two units still allows a discrete R-parity to survive, thereby eliminating the $d = 4$ operators. This occurs, for example, when the VEV of a 126 of SO(10) gives Majorana masses to ν^c. However, reps with dimension significantly above the adjoint's, such as SO(10)'s 126, are forbidden in stringy level-one models [3] and even seem unlikely to appear at any level [4]. Alternatively, products of VEVs for sneutrino-like fields, N^c and $\overline{N'}^c$, from vector-like pairs of 16_H, $\overline{16_H}$ Higgs fields can play the role of the 126 and give Majorana mass to the ν^c from the non-renormalizable superpotential term $16_i \cdot 16_i \cdot \overline{16_H} \cdot \overline{16_H}/M$. Simultaneously these fields can produce another SO(10)-allowed term, $16 \cdot 16 \cdot 16 \cdot 16_H/M$, which contains the dangerous $U^c D^c D^c \langle 16_H \rangle/M$ and $QLD \langle 16_H \rangle/M$ couplings. One should expect $\langle 16_H \rangle/M \sim \mathcal{O}(1/100)$, which would predict a proton lifetime $\sim 10^{-6}$ years.

The appearance of an additional $U(1)'_\psi$ from $E_6 \rightarrow$ SO(10)$\times U(1)'_\psi$ improves the situation somewhat, but still doesn't eliminating fast proton decay. A 27 of E_6 breaks into $(16_{Q_\psi=1} + 10_{Q_\psi=-2} + 1_{Q_\psi=4})$. The renormalizable η_i terms in (5) are forbidden by Q_ψ and $Q_T \equiv Q_\psi - (B-L)$ as well as by $(B-L)$. If required for ν^c masses, $\langle 16_H \rangle$ and $\langle \overline{16_H} \rangle$ break I_{3R}, $(B-L)$, and Q_ψ. Y and Q_T remain unbroken until the two MSSM Higgs in the $10_{Q_T=-2}$ or the singlet $1_{Q_T=4}$ (denoted as χ) take on a VEV. The lowest order non-renormalizable η_i terms are generated by

$$16 \cdot 16 \cdot 16 \frac{\langle N^c \rangle}{M} \frac{\langle 10_H \rangle \langle 10_H \rangle}{M^2}, \tag{6}$$

$$16 \cdot 16 \cdot 16 \frac{\langle N^c \rangle}{M} \frac{\langle \chi \rangle}{M}, \tag{7}$$

which if $\langle 10_H \rangle, \langle \chi \rangle \sim \mathcal{O}(1 \text{ TeV})$ yields $\eta_1 \eta_2 \ll 10^{-24}$.

Although the lowest order λ_j terms are forbidden by both Q_ψ and Q_T, the problem of triplet exchange exists. An effective mass term $\langle \chi \rangle H_3 H'_{3*} + \text{hc}$ for the Higgs triplets $H_3, H'_{3*} \in 10_H$ breaks Q_ψ and Q_T by four units. Exchange of these triplets would again induce dangerous effective $d = 5$ operators, with λ_j above acceptable limits for $\langle \chi \rangle \sim \mathcal{O}(1 \text{ TeV})$.

Alternatively, Pati argues that the extraodinary stability of the proton should emerge as a compelling feature of a model, resulting from symmetries endemic to the underlying theory. The extra, ubiquitous $U(1)$'s in heterotic string models can often form such symmetries. A typical example is the model presented in [5]. This (321) model contains a total of eight local Abelian symmetries beyond $U(1)_{I_{3R}}$ and $U(1)_{B-L}$. The first three of these, $U(1)_{i=1,2,3}$ are generational symmetries. Specifically, all i^{th} generation MSSM states carry an extra $Q_i = 1/2$ charge.

Both observable and hidden sector states carry $Q_{I_{3R}}$, Q_{B-L}, and $Q_{i=1 \text{ to } 6}$ charge, while only hidden sector states have additional $Q_{i=7,8}$ charge. In this basis, $U(1)_{i=1 \text{ to } 6}$ are all anomalous:

$$\text{Tr} U(1)_1 = \text{Tr} U(1)_2 = \text{Tr} U(1)_3 = 24 \tag{8}$$
$$\text{Tr} U(1)_4 = \text{Tr} U(1)_5 = \text{Tr} U(1)_6 = -12.$$

All of the anomaly can be rotated into in a single, unique

$$U(1)'_A \equiv c_A \sum_i \{\text{Tr} U(1)_i\} U(1)_i \tag{9}$$

$$= 2 [U(1)_1 + U(1)_2 + U(1)_3] \tag{10}$$
$$- [U(1)_4 + U(1)_5 + U(1)_6],$$

where c_A is a normalization coefficient. One choice of basis for the remaining non-anomalous components of the original set of $U(1)_i$ is

$$U(1)'_1 = U(1)_1 - U(1)_2, \tag{11}$$
$$U(1)'_2 = U(1)_4 - U(1)_5, \tag{12}$$
$$U(1)'_3 = U(1)_4 + U(1)_5 - 2U(1)_6, \tag{13}$$
$$\hat{U}(1)_\psi = U(1)_1 + U(1)_2 - 2U(1)_3, \tag{14}$$

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$$\hat{U}(1)_\chi \;=\; [U(1)_1 + U(1)_2 + U(1)_3]$$
$$+\; 2\,[U(1)_4 + U(1)_5 + U(1)_6]\,. \qquad (15)$$

While none of the new stringy charges can single-handedly prevent proton decay, Table 1 (borrowed from [1]) reveals the combination of either $(\hat{Q}_\psi, B-L)$ or $(\hat{Q}_\psi, \hat{Q}_\chi)$ can accomplish this. The charges of each pair are complimentary in their preventative roles: when one symmetry forbids a dangerous operator, the other allows it, and vice versa.

Constraints from satisfactory proton stability are minimal in this model: Either pair of symmetries need only survive from the string scale down to an intermediate scale significantly below M_{string}. The sneutrino-like fields can take on VEVs up to $\sim 10^{15\,\text{to}\,16}$ GeV, and hidden sector states can condense at a scale as high as $10^{15.5}$ GeV. The strongest assumption is that the suppression factor, $(\langle\phi\rangle/M)^n$, from singlet state VEVs is $\leq 10^{-9}$.

The potential of inducing dangerous $d = 5$ operators via exotic triplets in the infinite tower of Planck-mass states is entirely removed by $\hat{Q}_\psi$. A true symmetry in a string theory is respected by Planck-mass states as well. Thus, color-triplets in the massive tower can only appear in vector-like pairs, with opposite $\hat{Q}_\psi$-charges. Such mass-terms cannot induce $d = 5$ proton-decay operators, which violate $\hat{Q}_\psi$. Note that $\hat{Q}_\psi$ treats generations unequally, as shown in (14). This is in fact a characteristic necessary for $\hat{Q}_\psi$'s success. This feature distinguishes it from Q_ψ and implies that $\hat{Q}_\psi$-like symmetries cannot arise from within generation independent symmetries like E_6. There is a significant advantage of string theory over standard field theory regarding explanation of proton stability. That is, the three MSSM generations naturally arise within many classes of strings, while various string symmetries distinguish between these generations.

2.2. Local Discrete Symmetries

In addition to the standard proton decay operators appearing in the MSSM and the Higgs-related in a SUSY GUT, more exotic operators can appear in non-MSSMs unless forbidden by additional symmetries. States that could gener-ate such operators inevitably appear in realistic stringy (321) models. A typical stringy example is exotic vector-like pairs of $SU(3)$ triplet/anti-triplets, $(D^{'}, \bar{D}^{'})$. Such vector pairs can produce dangerous effective $d = 4$ proton-decay operators from superpotential terms of the form:

$$\{U^c D^c \bar{D}^{'},\, U^c \bar{D}^{'} \bar{D}^{'},\, QL\bar{D}^{'} \qquad (16)$$
$$U^c E^c D^{'},\, QQD^{'},\, QhD^{'}\} \;\;\times\;\; \left(\frac{\langle\phi\rangle}{M}\right)^n.$$

The first three terms in (16) are generalizations of those in (5), with $D^c \leftrightarrow \bar{D}^{'}$.

However, string models often have their own means for disallowing such operators. One such method of avoidance is local discrete symmetries, Δ [6–8]. Abelian Δ arise when a local $U(1)_{Z'}$, with charges quantized in units of q, is broken by a VEV of a state with charge Nq. Then a Z_N LDS still survives and $\sum_i Q_{Z',i}$ in superpotential terms must equal 0 (mod N). Faraggi illustrates the role of this type of symmetry in [7].[3] In the model of [14] discussed by Faraggi, a Z_4 LDS prevents, to *all* orders in n, the appearance of all terms in (16). The Z_4 LDS remains after a local $U(1)_{Z'}$ is broken by the VEV of a sneutrino carrying $Q_{Z'} = \frac{1}{2}$ charge. In [14] the $Q_{Z'}$ charges for all MSSM states are quantized in units of $n/2$, while charges for exotic triplets are quantized in units of $n'/4$, with $n, n' \in Z$. As a result, all of the potentially dangerous exotic terms carry $Q_{Z'}$ quantized in units of $(2n + 1)/4$. Thus, they are eliminated from the superpotential by the surviving Z_4 symmetry.

2.3. Anomalous $U(1)$'s

Anomalous $U(1)$'s such as those discussed in subsection 2.1 are typically present in compactified string models. After rotation of all anomalies into a single $U(1)'_A$ via (9), the Dine-Sieberg-Witten (DSW) mechanism [15] breaks $U(1)'_A$. However, this occurs at the expense of generating a Fayet-Iliopoulos (FI) D-term that must be cancelled by VEVs of states Φ_α carrying the anoma-

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[3]Discussion of the roles played by stringy non-Abelian LDS's will appear in [19].

lous charge,

$$D_{\mathrm{A}} = \sum_{\alpha} Q_{\mathrm{A}}^{\alpha} |\Phi_{\alpha}|^2 + \frac{g^2}{192\pi^2} \mathrm{e}^{\phi} \mathrm{Tr} Q_{\mathrm{A}} \qquad (17)$$

$$\to 0 . \qquad (18)$$

Additional D- and F-flatness constraints require

$$D_i = \sum_{\alpha} Q_{\mathrm{i}}^{\alpha} |\Phi_{\alpha}|^2 = 0, \qquad (19)$$

$$F_{\alpha} = \langle \partial W / \partial \Phi_{\alpha} \rangle = 0. \qquad (20)$$

The presence of an anomalous $U(1)$ has profound effects on a model. It results in (1) gauge rank reduction via the induced VEVs, (2) new non-renormalizable mass terms via the induced VEVs, (3) a new source to non-universality of soft-breaking scalar masses, and (4) new contributions to the cosmological constant through the FI D-term. Additionally, many past models have used anomalous $U(1)$'s as generation symmetries, while recently an anomalous $U(1)$ has been used to construct models with SUSY-breaking [16,17].

Until last year, determining if a given model contained a $U(1)'_A$ meant computing the charge trace for each $U(1)_i$. As a first step towards classification of stringy anomalous $U(1)'_A$, Kobayashi and Nakano [8] find constraints on the appearance of a $U(1)'_A$ in orbifold models of $\mathrm{E}^8 \times \mathrm{E}^8$ class. They show that a $U(1)'_A$ will be absent if there is no massless twisted state carrying both observable and hidden sector charges or if certain types of discrete symmetries are found. Moreover, a given $U(1)_i$ in the observable (hidden) sector is anomaly free if there is a hidden (observable) group G such that all massless G-charged state have vanishing charge Q_i. Relatedly, no model with an unbroken E_8 can have an anomalous $U(1)'_A$.

Another constraint is if a model is invariant under a Z_2 momentum duality $\bar{P}_i \to -\bar{P}_i$ in the compactified direction X_i, then the charges of the corresponding $U(1)_i$ sum to zero (with the exception that if $\bar{P}_i$ is a spinorial momentum, then a second, simultaneous, $\bar{P}_j \to -\bar{P}_j$ is needed).

As a stepping stone for a general $U(1)'_A$ classification procedure for orbifolds, Kobayashi and Nakano demonstrate that analysis of orbifolds in a Wilson-line background can essentially be reduced to analysis in the absence of a Wilson-line.

Classification of $U(1)'_A$ in terms of modular invariant shifts in a model is introduced. Of the five distinct classes of shifts found, some classes are shown to lead to anomalies while others do not. Lastly, it is suggested that LDS's will generally remain after breaking of the $U(1)'_A$ and that these LDS's apply additional constraints on flat directions. There may also be a relationship between the appearance of discrete R-symmetries and $U(1)'_A$.

Classification of $U(1)'_A$ in free fermionic models is being studied by G.C. and Alon Faraggi. Translation of orbifold findings into constraints on basis vectors and phase coefficients (equivalently, k_{ij}-matrix entries) is underway.

2.4. Kinetic Mixing Among $U(1)_i$

Recently a new and possibly very general constraint on phenomenologically viable string models appeared [9]. Supersymmetry breaking hidden sectors containing local $\tilde{U}(1)_k$ hidden gauge factors could destabilize the observable gauge hierarchy under certain conditions. For both SUGRA- and GUT-mediated SUSY-breaking, kinetic mixing has the potential to raise the electroweak scale up to the hidden-sector SUSY-scale, $M_{\mathrm{EW}} \to M_{\mathrm{susy}}^{\mathrm{hid}}$. This can occur when a string model contains states that simultaneously carry both observable $U(1)_i$ (including $U(1)_Y$) and hidden sector $\tilde{U}(1)_k$ charges. This occurrence is usually the case, as discussions in prior subsections should suggest.

Such a model may contain in its pure gauge part of the Lagrangian a $U(1)$ "kinetic mixing term," with mixing strength measured by the dimensionless coefficient χ:

$$\begin{aligned} \mathcal{L}_{\mathrm{gauge}} &= -\frac{1}{4} F_{(i)}^{\mu\nu} F_{(i)\mu\nu} - \frac{1}{4} F_{(k)}^{\mu\nu} F_{(k)\mu\nu} \\ &+ \frac{\chi}{2} F_{(i)}^{\mu\nu} F_{(k)\mu\nu}, \end{aligned} \qquad (21)$$

Supersymmetrization corresponds to,

$$\mathcal{L}_{\mathrm{gauge}} = \frac{1}{32} \int d^2\theta (W_i W_i + W_k W_k \\ - 2\chi W_i W_k), \qquad (22)$$

where $W_i \equiv \bar{D}^2 D V_i$ is the chiral gauge field strength formed from the vector superfield V_i.

Assuming that the hidden sector $\tilde{U}(1)_k$ breaks when a hidden sector state with charge $\tilde{Q}_k$ takes a VEV, the pure gauge term in the Lagrangian may be brought to canonical form by performing the shift, $V_i^\mu \rightarrow V_i^{'\mu} = V_i^\mu - \chi V_k^\mu$. This shift leads to new interactions between the observable and hidden sectors:

(1) observable states φ get hidden sector charges proportional to their observable sector charges and couplings,

$$\mathcal{D}_k^\mu = \partial^\mu + i(g_k \tilde{Q}_k + g_i \chi Q_i) V_k^\mu, \qquad (23)$$

(2) observable states couple to hidden sector gauginos λ_k,

$$\mathcal{L} = i\sqrt{2} g_i Q_i \chi \varphi^\dagger \tilde{\varphi} \lambda_k + \text{ h.c.}, \qquad (24)$$

(3) the hidden sector D_k term is shifted,

$$D_i' = -g_i \sum_a Q_{i_a} |\varphi_a|^2, \qquad (25)$$

$$D_k = -g_k \sum_b \tilde{Q}_{k_b} |\Phi_b|^2 + \chi D_i', \qquad (26)$$

(4) the hidden sector D_k term makes a contribution to soft scalar masses (m_a^2) by an amount,

$$(m_a^2)_{KM} = g_i \chi Q_{i_a} \langle D_k \rangle. \qquad (27)$$

Whether or not these mixing effects are significant depends on the ratio $\eta \equiv \chi \langle D_k \rangle / M_Z^2$. $\eta \gg 1$ suggests kinetic mixing could be the dominant messenger of SUSY-breaking and produce a hierarchy destabilization as $M_{\text{scalar}} \gg M_Z$. While $\eta \sim \mathcal{O}(1)$ does not destablize the gauge group hierarchy, it still yields significant corrections to soft scalar masses. On the other hand, if $\eta \ll 1$ then all mixing effects are ignorable.

Stability of a model against KM effects is assured if hidden sector states appear as conjugate pairs for each $\tilde{U}(1)_k$ involved in hidden sector SUSY-breaking. This produces a complete cancellation of kinetic mixing terms at all orders, (i.e., $\chi = 0$). Existence of such conjugate pairs implies a certain hidden sector Z_2 LDS $\tilde{\Delta}_l$. For generic string models lacking such conjugate states, it is estimated that even if $\chi = 0$ at M_{string} at tree level, one-loop effects should generate non-zero values at lower scales,

$$\frac{\chi}{g_i g_k} = \frac{b_{ik}}{16\pi^2} \ln \frac{M_{\text{string}}^2}{\mu^2} + \frac{1}{16\pi^2} \Delta_{ik}. \qquad (28)$$

The first term on the right-hand side of eq. (28) is the massless sector contribution, where

$$b_{ik} = -\text{Str}_{\text{massless}} \bar{Q}_H^2 Q_i \tilde{Q}_k, \qquad (29)$$

with $\bar{Q}_H^2$ the spacetime helicity operator. The second term is an effect of the infinite tower of massive string states. Typical four-dimensional string models yield

$$10^{-1} \lesssim |\Delta_{ik}| \lesssim 10^{+1}, \qquad (30)$$

which corresponds to

$$3 \times 10^{-4} \lesssim \frac{\chi(M_Z)}{g_b(M_Z)} \lesssim 3 \times 10^{-2}. \qquad (31)$$

Thus, seemingly one should expect $\eta(M_Z) \gg 1$ in generic string models following either SUGRA or gauge-mediated SUSY-breaking paths.

With respect to this estimated range of typical χ-values, stringy three-generation (321) models (and GUT/semi-GUT generalizations) with $N = 1$ SUSY appear in striking contrast to the average model. Amazingly, in generic $N = 1$ (321) models, contributions to χ cancel level-by-level, starting with the massless states. However, the authors of [9] argue that, while cancellation is exact for tree-level masses, cancellation is unstable against mass splittings from GUT-breaking, non-zero VEVs related to the DSW mechanism, low-energy SUSY-breaking, and resulting RGE-flows. χ/g_k contributions from mass splittings resulting from these first four effects are estimated to be of $\mathcal{O}(10^{-3})$, $\mathcal{O}(10^{-8}$ to $10^{-5})$, $\mathcal{O}(10^{-8})$, and $\mathcal{O}(10^{-15})$, respectively. Much of these range of values would yield either η in the dangerous $\gg 1$ range or in the phenomenologically interesting $\eta \sim 1$ range. Thus, avoiding potentially dangerous kinetic mixing effects in quasi-realistic heterotic string models may imply the appearance of appropriate Z_2 LDS's in such models. However, further research into kinetic mixing appears necessary before definitive statements can be made regarding the effective low energy values of η for the (321) models.

3. Progress in Actual Model Building

While definitions of string model building can vary, the practice can be grouped into two primary classes: string *derived* models and string

inspired models. A string *derived* model essentially involves a top-down approach, starting with a true string model that yields an effective (321), semi-GUT, or GUT model in the low energy limit. On the other hand, a string *inspired* (321), semi-GUT, or GUT model follow a more bottom-up approach. It is a field theoretic model upon which certain string-like constraints are imposed. Generally these constraints include limits on the dimensions of gauge group reps[4] and resolution of the apparent disparity between the string scale $M_{\text{string}} \approx g_s \times 5.5 \times 10^{17}$ GeV and the MSSM unification scale $M_{\text{unif}} \equiv 2.5 \times 10^{16}$ GeV.

Of the several proposed solutions to the scale disparity, findings of the last few years [10] suggest stringy GUTs [11] and/or non–MSSM states between 1 TeV and M_{unif} are perhaps the only feasible ones on the list (neglecting unknown non-perturbative effects). The Non-MSSM approach often generates semi-GUT unification near M_{unif} or additional gauge factors above M_{unif}.

The following subsections present a cursory review of developments along both the derived and inspired lines of model construction. Stringy GUTs of the last year generally fall into the derived class [18,19], whereas semi-GUTs [21,22,20] and (Near/Extended-)MSSM [23,25] tend to be of the more inspired nature.

3.1. String Derived GUT Models

While several groups in recent years have been pursuing string-derived SUSY GUTs [26], a sticking point has been the difficulty producing three net generations in conjunction with adjoint (or higher dimension) Higgs. Success was finally realized last year by Kakushadze and Tye using Z_6 asymmetric orbifolds that include a Z_3 outer-automorphism [18]. The latter produces level-

three current algebras yielding three-generation $SU(5)$, $SU(6)$, $SO(10)$, and E_6 GUTs. All models in this class contain a single adjoint (and no higher dimension reps). Discovering these models is undoubtedly the most significant string GUTs advancement of the past year. Zurab Kakushadze will speak next about these models in detail.

My study of level-two and level-four[5] free fermionic $SO(10)$ models is ongoing. I have found no such models with three generations yet. While general arguments indicate three generations is more difficult to produce at these levels than at level-three, a no-go theorem does not exist. One may eventually result from my search though. However, if level-four models with three generations are ever found, they may be able to contain more than one adjoint Higgs, possibly in conjunction with 54's [4].

3.2. String Inspired Semi-GUT Models

Four closely related SUSY semi-GUT models were presented recently as solutions to the MSSM-vs-string unification scale disparity. Rizos et. al [20] have developed $SU(6) \times SU(2)_l$ and $SU(6) \times SU(2)_r$ models, while Despande et. al [21] and Murayama [22] start at the string scale one step closer to the MSSM, with their $SU(4)_c \times SU(2)_l \times SU(2)_r$ models.

The states in these models are consistent with level-one algebras. In Rizos et. al, the gauge couplings of $SU(6)$ and $SU(2)_{l\,(r)}$ meet at the string scale. $SU(6) \times SU(2)_{l\,(r)}$ breaks to $SU(4) \times SU(2)_l \times SU(2)_r$ at a scale $M_{\text{X}} \approx M_{\text{unif}}$ when $(15,1)$ and $(\overline{15},1)$ reps acquires VEVs along $SU(4)_c \times SU(2)_l \times SU(2)_r$ singlet directions. At a slightly lower scale $M_{\text{R}} \sim 10^{14.6\,\text{to}\,16.3}$ GeV, $SU(4)_c \times SU(2)_l \times SU(2)_r$ breaks to $SU(3)_c \times SU(2)_l \times U(1)_Y$ as $(4,1,2)$ and $(\overline{4},1,2)$ reps acquire VEVs along the $SU(3)$ singlet direction. Below M_{R}, the only non-MSSM states are a pair of exotic triplets with fractional electric charge. These become massive at $10^{5\,\text{or}\,6}$ GeV.

In the limit of $M_{\text{X}} \to M_{\text{string}}$, one of the few differences between the $SU(6) \times SU(2)_l$ model and the $SU(4) \times SU(2)_l \times SU(2)_r$ models is the lowering of M_{R} to $10^{11\,\text{to}\,12}$ GeV for the latter case.

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[4] For massless stringy matter to appear as a rep R of a gauge group, unitarity constraints require $\lambda^R \cdot \Psi \leq K$, where λ^R is the highest weight of R, Ψ is the highest root in the gauge group, and K is the level of the gauge group. For example only the 1, 10, 16, and $\overline{16}$ reps can appear for $SO(10)$ at level-one. Adjoint Higgs (and higher dimension reps) can only appear at $K \geq 2$ for any gauge group. Further, the mass of a state is a linear function of the quadratic Casimir of the rep, while it is an inverse function of the level-K. Thus, while all SO(10) reps up through the 210 are allowed by unitarity at level-two, only the singlet through the 54 can be massless.

[5] $SO(10)$ level-three does not have a free fermionic realization.

Murayama claims these models are the most efficient way of resolving the scale disparity. Each exotic state in these models plays a critical role besides simply increasing the unification scale to near M_{string}.

3.3. String Inspired Extended-MSSMs

Allanach et. al [23] have developed a class of extended MSSM models [24] that raises the MSSM coupling unification scale to approximately M_{string} by the addition of extra vector-like exotic MSSM states. Based on typical stringy anomaly cancellation requirements, the mass scale of the exotics, M_{I}, is assumed comparable to M_{unif}. The exotics act on the MSSM running couplings like an RGE lens located at M_{I}. This lens refocuses the gauge coupling unification point to a value M_{X} near M_{string}. In these models, the (321) symmetry is enhanced above M_{I} by an extra $U(1)_X$ generational symmetry, whose coupling merges with the others at M_X.

The range of MSSM vector-like exotics considered is limited to those permitted by level-one algebras. These are $[(3,2)_{Y,X},(\overline{3},2)_{-Y,-X}]$ quark-like pairs, $[(3,1)_{Y,X},(\overline{3},1)_{-Y,-X}]$ anti-quark-like pairs, $[(1,2)_{Y,X},(1,2)_{-Y,-X}]$ lepton doublet/ Higgs-like pairs, and $[(1,1)_{Y,X},(1,1)_{-Y,-X}]$ singlet pairs. Let n_Q, n_q, n_2, and n_1 represent the number of exotics in each of these classes, respectively. Several different combinations of $\{n_Q, n_q, n_H\}$ yield $M_{\text{X}} \approx M_{\text{string}}$, with the most general requirement being $n_q > n_Q + n_2$. The simplest combination giving $M_X \approx M_{\text{string}}$ is $n_q = 3$, $n_Q = 2$, $n_2 = 0$ plus one vector-like singlet, $(\Theta_{X=1}, \overline{\Theta}_{X=-1})$, to break $U(1)_X$. Acceptable Yukawa mass textures result in more complicated models containing several vector-like Higgs (with a compensating increase in vector-like anti-quarks).

Assuming a higher level for $SU(3)_c$ and $SU(2)_l$ offers an alternative answer for the scale disparity issue. Bachas et. al and Bastero-Gil et. al [25] show that (321) unification can be pushed up to M_{string} by the addition of just $SU(3)_c$ and $SU(2)_l$ adjoint scalars with mass around 10^{14} GeV. Since this mass scale is comparable to the scale at which coupling becomes non-perturbative for gauge groups $SU(3)$ or $SU(5)$, it is suggested

that the masses of these adjoints could be related to hidden sector gaugino condensation.

3.4. μ-Term and Mass Hierarchy

General classes of string inspired solutions have also been proposed for the μ-term problem [27–29] and mass heirarchy [30,31,27,32,33]. One of the primary inspired aspects of this pursuit is the general appearance and critical role played by extra $U(1)$'s. Later this morning J.R. Espinosa will speak about related work here at Penn [27]. So I leave further discussion of this topic to him.

4. Concluding Comments

"Old fashioned" heterotic string model building is definitely in its maturation stage. It has entered an era involving finely detailed (re-)analysis of (past) models. Understanding phenomenological patterns within model classes has received new attention, especially so regarding the roles of extra $U(1)$'s and local discrete symmetries (LDS's) in the models. Extra local Abelian symmetries and LDS's play vital roles in preventing fast proton decay and flavor changing neutral currents. Moreover, whenever any extra $U(1)$'s are anomalous, the phenomenology of the model is drastically altered. This has inspired recent study of conditions required for anomalous $U(1)$'s to appear. Another avenue of model analysis has surfaced with the discovery of possible kinetic mixing between observable and hidden sector $U(1)$'s . If this mixing is strong enough, possibly destablization of the observable sector gauge hierarchy could occur. LDS's in the hidden sector offer a means of eliminating this danger.

As a whole, string model building in the weakly coupled realm continues to play an important role in high energy theory. Even in the the context of the ongoing second string revolution, there is much to be gained from both the top-down and bottom-up stringy approaches to model building.

5. ACKNOWLEDGMENTS

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Operators	Generation Combinations	Y	$B-L$	$\hat{Q}_\psi$	$\hat{Q}_\chi$	$\hat{Q}_\chi + \hat{Q}_\psi$	If Allowed
$U^c D^c D^c, QLD^c, LLE^c$	(a) All except (b)	√	×	×	×	×	unsafe
	(b) 3 fields from 3 different families	√	×	√	×	×	unsafe
$(U^c D^c D^c$ or $QLD^c)(N^c/M)$	All	√	√	×	√	×	unsafe
$(U^c D^c D^c$ or $QLD^c)(N^c/M)\times$ $[(\overline{h_1}/M)^2$ or $(\phi/M)^n]$	All	√	√	√	×	×	safe
$(U^c D^c D^c$ or $QLD^c)(N^c/M)\times$ $(T_a\overline{T}_b/M^2)^2$	Some(†)	√	√	√	√	√	safe
$QQQL/M$	All	√	√	×	√	×	unsafe
$(QQQL/M)(\overline{N'}^c_\alpha/M)_{\alpha=1,2}$	e.g.$(1,2,1,3)$	√	×	√	×	×	unsafe
$(QQQL/M)(\overline{N'}^c_\alpha/M)(N^c_\beta/M)$	All	√	√	×	√	×	safe(?)
$(QQQL/M)(T_a\overline{T}_b/M^2)^2$	Some(†)	√	√	√	√	√	safe
$U^c U^c D^c \overline{E}/M$	All	√	√	×	×	(∗)	unsafe
$LL\overline{h_i}\,\overline{h_i}/M$	All	√	×	×		(∗)	safe

Table 1

(Borrowed from [1].) The roles of Y, $B-L$, $\hat{Q}_\psi$, $\hat{Q}_\chi$ and $\hat{Q}_\chi + \hat{Q}_\psi$ in allowing (√) or forbidding (×) the relevant (B,L) violating operators. The mark † signifies that the corresponding operator is allowed if either two of the four MSSM fields are in generation 1 or 2 and two fields are in generation 3, with $a = 1$ and $b = 3$ for the hidden sector states T_a and $\overline{T}_b$; or all four fields are in generation 1 or 2 with $a = 1$ and $b = 2$. The mark (∗) signifies that $(\hat{Q}_\chi + \hat{Q}_\psi)$ forbids $U^c U^c D^c E^c/M$ for all generation combinations except when all four fields belong to generation 3, and that it forbids $LL\overline{h_i}\,\overline{h_i}$ in some generation combinations, but not in others. The Higgs doublet $\overline{h_i}$ carries generation charge Q_i.